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Higgs and Z boson decays into light gluinos

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Abstract

We calculate the decay rate of scalar and pseudoscalar Higgs bosons into a pair of gluinos, within the Minimal Supersymmetric Standard Model. In the theoretically and experimentally allowed light gluino window, $m_{\tilde{g}} \sim 3$–5 GeV, gluino pairs can completely dominate the decays of the light scalar Higgs boson and play a prominent role in the decay of the pseudoscalar Higgs boson. This would alter the limits obtained from Z decays on the lightest CP–even and CP–odd Higgs bosons, and could jeopardize the search for these Higgs particles at future hadron colliders. In contrast, the branching ratio for the two–body decay of Z bosons into pairs of light gluinos is less than 0.1%.

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1. Introduction

The last couple of years have seen renewed interest [1] in the possible existence of very light gluinos, with mass $m_{\tilde{g}} \leq 5$ GeV. A recent analysis by Farrar [2] concludes that only the upper end of this range is still open, roughly 3 GeV $\leq m_{\tilde{g}} \leq 5$ GeV. However, even this window would be sufficient to allow for a substantial modification of the running of α_S, which was the original motivation [1] for this scenario. Moreover, since the decays of light gluinos might only produce a small amount of missing energy [3], bounds on squark masses derived from searches [4] for events with large missing E_T at $p\overline{p}$ colliders are not valid in this scenario.

In this paper we point out that the existence of light gluinos could also substantially complicate the search for the Higgs bosons predicted by supersymmetric models. To be specific, we work within the framework of the minimal supersymmetric standard model or MSSM, which contains two $SU(2)$ doublets of Higgs superfields [5, 6]. We will not assume unification of gaugino masses, which would impose severe constraints on models with light gluinos [7], nor do we require radiative breaking of the $SU(2)$ symmetry, which appears to be inconsistent with the existence of light gluinos, at least in the minimal model [8]. We find that the $\tilde{g}\tilde{g}$ final state can completely dominate the decay of the light scalar Higgs boson; in some cases the branching ratios for the usually dominant $b\overline{b}$ and $\tau^+\tau^-$ final states are reduced by a factor of 1,000 or more. Light gluinos might also be produced in 70% of all decays of the pseudoscalar Higgs boson of the MSSM. This has obvious consequences for Higgs searches that rely on b or τ tagging. The decays of supersymmetric Higgs bosons to gluino pairs have been discussed in the past in ref. [9]. However, the possibility that the gluinos are light and the fact that the mixing between the two supersymmetric scalar partners of the top quark (which, as we will see later, is the most important feature in this decay) might be large have not been considered. Therefore the Higgs $\rightarrow \tilde{g}\tilde{g}$ branching ratios obtained in ref. [9] were very small.

It has been known for some time [10] that light gluinos can also be produced in two–body decays of Z bosons; this decay has recently been studied in ref. [11]. We re–compute the corresponding partial width, including the effect of mixing between $SU(2)$ doublet and singlet squarks. Our numerical results agree with those of ref. [11], but our final expression is much more compact. We find a maximal $Br(Z \rightarrow \tilde{g}\tilde{g}) \simeq 6 \cdot 10^{-4}$ within our assumptions; such a small branching ratio can only be probed by a dedicated search for $\tilde{g}\tilde{g}$ final states. Unfortunately the failure to detect $Z \rightarrow \tilde{g}\tilde{g}$ decays would not exclude the possibility that gluino pairs contribute significantly or dominantly to Higgs boson decays.

The remainder of this article is organized as follows. In sec. 2 expressions for the partial widths for the decays of Higgs and Z bosons into pairs of light (effectively massless) gluinos are given. In sec. 3 we present numerical results for the separate branching ratios and discuss possible correlations between them. Sec. 4 contains a brief summary and conclusions. The Appendix collects expressions for the decays of Higgs bosons into massive gluinos, as well as the necessary two– and three–point functions.

2. Formalism

The diagrams contributing to the one–loop induced decays of Higgs and Z bosons into a pair of gluinos are shown in Fig. 1. Since the gluinos are of Majorana nature and hence identical fermions, one has to antisymmetrize the decay amplitude; this is achieved by adding the contributions of the two diagrams where the momenta of the gluinos are interchanged and by multiplying these contributions by an overall minus sign. In the case of Higgs bosons, this reduces to just multiplying the amplitudes of the diagrams shown in Fig. 1 by a factor of two; for the Z boson one needs in addition to discard the terms that are not proportional to γ_5.

In Fig. 1 one has to include the contributions of both squarks of a given flavor. As well known [12], the supersymmetric partners of left– and right–handed massive quarks mix. The mass eigenstates $\tilde{q}_1$ and $\tilde{q}_2$ are related to the current eigenstates $\tilde{q}_L$ and $\tilde{q}_R$ by

$$\tilde{q}_1 = \tilde{q}_L \cos\theta_q + \tilde{q}_R \sin\theta_q \quad , \qquad \tilde{q}_2 = -\tilde{q}_L \sin\theta_q + \tilde{q}_R \cos\theta_q . \tag{1}$$

The mixing angle θ_q as well as the masses $m_{\tilde{q}_1}$, $m_{\tilde{q}_2}$ of the physical squarks can be calculated from the following mass matrices[1], written in the convention of ref. [13]:

$$\mathcal{M}^2_{\tilde{t}} = \begin{pmatrix} m^2_{\tilde{t}_L} + m_t^2 + 0.35 D_Z & -m_t(A_t + \mu\cot\beta) \\ -m_t(A_t + \mu\cot\beta) & m^2_{\tilde{t}_R} + m_t^2 + 0.16 D_Z \end{pmatrix}; \tag{2a}$$

$$\mathcal{M}^2_{\tilde{b}} = \begin{pmatrix} m^2_{\tilde{t}_L} + m_b^2 - 0.42 D_Z & -m_b(A_b + \mu\tan\beta) \\ -m_b(A_b + \mu\tan\beta) & m^2_{\tilde{b}_R} + m_b^2 - 0.08 D_Z \end{pmatrix}, \tag{2b}$$

where $D_Z = M_Z^2 \cos 2\beta$, $\tan\beta$ being the ratio of the vacuum expectation values of the two neutral Higgs fields of the MSSM [5]. $m_{\tilde{t}_L, \tilde{t}_R, \tilde{b}_R}$ are soft breaking masses, $A_{b,t}$ are parameters describing the strength of nonsupersymmetric trilinear scalar interactions, and μ is the supersymmetric Higgs(ino) mass, which also enters trilinear scalar vertices. Notice that the off–diagonal elements of these squark mass matrices are proportional to the quark mass. In the case of the supersymmetric partners of the light quarks mixing between the current eigenstates can therefore be neglected. However, mixing between $\tilde{t}$ squarks can be sizable and allows one of the mass eigenstates to be much lighter than the top quark. Sbottom mixing can also be significant if $\tan\beta \gg 1$; even in supergravity models with radiative symmetry breaking $\tan\beta$ can be as large as m_t/m_b [13].

In the presence of squark mixing, the squark-quark-gluino interaction Lagrangian is given by

$$\mathcal{L}_{\tilde{g}\tilde{q}q} = -i\sqrt{2} g_s T^a \overline{q} \left[(\cos\theta_q \tilde{q}_1 - \sin\theta_q \tilde{q}_2) \frac{1+\gamma_5}{2} - (\sin\theta_q \tilde{q}_1 + \cos\theta_q \tilde{q}_2) \frac{1-\gamma_5}{2} \right] \tilde{g}_a + \text{h.c.}, \tag{3}$$

where g_s is the strong coupling constant and T^a are $SU(3)_C$ generators. Note that in eq. (3) we have assumed $M_{\tilde{g}} > 0$.

[1] We ignore generation mixing between squarks, which in case of the MSSM is only induced radiatively by weak interactions.

2a. Higgs boson decays into light gluinos

Summing over colors and taking into account the fact that there are two identical particles in the final state, the partial decay widths of the CP–even Higgs bosons $H_1, H_2 \equiv h$ and the CP–odd boson $H_3 \equiv P$ into a pair of gluinos in the limit $m_{\tilde{g}} \to 0$ are given by

$$\Gamma(H_k \to \tilde{g}\tilde{g}) = \frac{1}{2}\alpha M_{H_k} \left(\frac{\alpha_S}{\pi}\right)^2 \left(\sum_{q=t,b} A_k^q\right)^2 . \tag{4}$$

The amplitudes A_k^q can be written as

$$A_{1,2}^q = \frac{1}{2}(s_{1,2})^q \sin 2\theta_q \left[(m_q^2 + m_{\tilde{q}_2}^2) C_0(m_q, m_q, m_{\tilde{q}_2}) - (m_q^2 + m_{\tilde{q}_1}^2) C_0(m_q, m_q, m_{\tilde{q}_1})\right] - \frac{1}{2} m_q \sin 2\theta_q \\ \times \left[\left(\tilde{s}_{1,2}^{11}\right)^q C_0(m_{\tilde{q}_1}, m_{\tilde{q}_1}, m_q) - \left(\tilde{s}_{1,2}^{22}\right)^q C_0(m_{\tilde{q}_2}, m_{\tilde{q}_2}, m_q) + 2\left(\tilde{s}_{1,2}^{12}\right)^q \cot 2\theta_q C_0(m_{\tilde{q}_1}, m_{\tilde{q}_2}, m_q)\right];$$

$$A_3^q = \frac{1}{2}(s_3)^q \sin 2\theta_q \left[(m_q^2 - m_{\tilde{q}_2}^2) C_0(m_q, m_q, m_{\tilde{q}_2}) - (m_q^2 - m_{\tilde{q}_1}^2) C_0(m_q, m_q, m_{\tilde{q}_1})\right] \\ + \left(\tilde{s}_3^{12}\right)^q m_q C_0(m_{\tilde{q}_1}, m_{\tilde{q}_2}, m_q), \tag{5}$$

where the scalar function $C_0(m_1, m_2, m_3)$ is given in integral form by

$$C_0(m_1, m_2, m_3) = -\int_0^1 dy \int_0^y dx \left[M_{H_k}^2 x(x-y) + (m_2^2 - m_3^2)y + (m_1^2 - m_2^2)x + m_3^2 - i\epsilon\right]^{-1} \tag{6}$$

A complete expression after integration over the Feynman variables is given in the Appendix. In the limit $M_{H_k}^2 \ll m_t^2$ this function simplifies to

$$C_0(m_1, m_2, m_3) = \frac{1}{m_2^2 - m_1^2}\left[\frac{m_1^2 \log m_1^2 - m_3^2 \log m_3^2}{m_1^2 - m_3^2} - \frac{m_2^2 \log m_2^2 - m_3^2 \log m_3^2}{m_2^2 - m_3^2}\right], \tag{7}$$

and for $m_1 = m_2$

$$C_0(m, m, m_3) = \frac{1}{m_3^2 - m^2} + \frac{m_3^2}{(m_3^2 - m^2)^2} \log \frac{m^2}{m_3^2} \tag{8}$$

The $H_k qq$ couplings $(s_k)^q$ are given by [5] ($s_W^2 = 1 - c_W^2 = \sin^2\theta_W$)

$$(s_k)^q = \frac{m_q}{2 s_W M_W} r_k^q, \tag{9}$$

with

$$r_1^t = \frac{\sin\alpha}{\sin\beta} \;,\; r_2^t = \frac{\cos\alpha}{\sin\beta} \;,\; r_3^t = \frac{1}{\tan\beta}; \\ r_1^b = \frac{\cos\alpha}{\cos\beta} \;,\; r_2^b = -\frac{\sin\alpha}{\cos\beta} \;,\; r_3^b = \tan\beta. \tag{10}$$

The $H_k\tilde{q}_i\tilde{q}_j$ couplings $(\tilde{s}_k^{ij})^q$ are [14]

$$\begin{aligned}
(\tilde{s}_1^{11})^q &= \frac{M_Z\cos(\alpha+\beta)}{s_W c_W}\left(I_{3q}^L\cos^2\theta_q - e_q s_W^2\cos 2\theta_q\right) + \frac{m_q^2 r_1^q}{s_W M_W} - \frac{m_q\sin 2\theta}{2s_W M_W}(r_1^q A_q + r_2^q\mu);\\
(\tilde{s}_1^{22})^q &= \frac{M_Z\cos(\alpha+\beta)}{s_W c_W}\left(I_{3q}^L\sin^2\theta_q + e_q s_W^2\cos 2\theta_q\right) + \frac{m_q^2 r_1^q}{s_W M_W} + \frac{m_q\sin 2\theta}{2s_W M_W}(r_1^q A_q + r_2^q\mu);\\
(\tilde{s}_1^{12})^q &= \frac{M_Z\cos(\alpha+\beta)}{s_W c_W}\sin 2\theta_q\left(e_q s_W^2 - \frac{1}{2}I_{3q}^L\right) - \frac{m_q}{2s_W M_W}(r_1^q A_q + r_2^q\mu)\cos 2\theta_q; \qquad (11)
\end{aligned}$$

$$(\tilde{s}_2^{ij})^q = (\tilde{s}_1^{ij})^q\ [\sin\alpha \to +\cos\alpha\ , \cos\alpha \to -\sin\alpha]; \qquad (12)$$

$$(\tilde{s}_3^{11})^q = (\tilde{s}_3^{22})^q = 0 \quad , \ (\tilde{s}_3^{12})^q = \frac{m_q}{2s_W M_W}(\mu - r_3^q A_q). \qquad (13)$$

A few remarks are in order:

(i) As already mentioned, eqs.(5)–(8) are valid only if the gluinos are nearly massless; this is an excellent approximation for the case of interest, $m_{\tilde{g}} \leq 5$ GeV $\ll m_t, m_{H_k}, m_{\tilde{q}}$. Complete expressions for a finite gluino mass are given in the Appendix. Note that the contribution of diagram Fig. 1b is ultraviolet finite but the contribution of diagram Fig. 1a is finite only once the contributions of both squarks of a given flavor have been added.

(ii) The main conventional decay mode of Higgs particles with masses below 130 GeV is the decay into $b\bar{b}$ pairs; in the limit $M_{H_k} \gg m_b$, the corresponding decay width is given by

$$\Gamma(H_k \to b\bar{b}) = \frac{3}{2}\alpha M_{H_k}\left(\frac{m_b}{2s_W M_W}\right)^2 (r_k^b)^2. \qquad (14)$$

Note that one has to include the QCD corrections to this decay width. The bulk of these corrections can be absorbed [15] into running quark masses evaluated at the scale $\mu = M_{H_k}$. For Higgs masses arround 100 GeV, the b–quark mass $m_b(m_b) = 4.5$ GeV will drop to the effective value $m_b(M_{H_k}) \simeq 3.2$ GeV; this results in a decrease of the decay width by approximately a factor of two.[2]

(iii) In the limit where either A_q and μ or m_q are set to zero, there is no mixing between left– and right–handed squarks. In this case, the amplitudes A_k^q are $\propto m_{\tilde{g}}$ and hence very small; see also ref. [9]. Therefore only the contribution of the top quark and its SUSY partners have to be taken into account; (s)bottom loop contributions are sizable only if tan$\beta \gg 1$.

2b. The decay $Z \to \tilde{g}\tilde{g}$

The partial decay width of the Z boson into a pair of massless gluinos is given by

$$\Gamma(Z \to \tilde{g}\tilde{g}) = \frac{\alpha M_Z}{48 c_W^2 s_W^2}\left(\frac{\alpha_S}{\pi}\right)^2\left[\sum_q\left(\sum_{i=1}^2 B_i^q + \sum_{i,j=1}^2 \tilde{B}_{ij}^q\right)\right]^2, \qquad (15)$$

[2]In contrast to the running of α_S, light gluinos have very little effect on the running of m_b.

where [here $s = M_Z^2$]

$$
\begin{aligned}
B_i^q &= (v_i^2 + a_i^2) a_q \left[2C_2(s, m_q, m_q, m_{\tilde{q}_i}) - B_0(s, m_q, m_q) - (m_q^2 + m_{\tilde{q}_i}{}^2) C_0(s, m_q, m_q, m_{\tilde{q}_i}) \right] \\
&\quad + 2(v_i a_i) v_q \left[2C_2(s, m_q, m_q, m_{\tilde{q}_i}) - B_0(s, m_q, m_q) + (m_q^2 - m_{\tilde{q}_i}{}^2) C_0(s, m_q, m_q, m_{\tilde{q}_i}) \right] ; \\
\tilde{B}_{ij}^q &= -2a_{ij}(v_i a_j + v_j a_i) C_2(s, m_{\tilde{q}_i}, m_{\tilde{q}_j}, m_q),
\end{aligned} \tag{16}
$$

with

$$
a_{11} = 2(2I_q^{3L} \cos^2\theta_q - 2s_W^2 e_q) \ , \ a_{22} = 2(2I_q^{3L} \sin^2\theta_q - 2s_W^2 e_q) \ , \ a_{12} = -2I_q^{3L} \sin 2\theta_q, \tag{17}
$$

$$
v_q = 2I_q^{3L} - 4s_W^2 e_q \quad , \quad a_q = 2I_q^{3L}, \tag{18}
$$

and

$$
v_1 = \frac{1}{2}(\cos\theta_q - \sin\theta_q) = a_2 \quad , \quad a_1 = \frac{1}{2}(\cos\theta_q + \sin\theta_q) = -v_2. \tag{19}
$$

Alternatively, the result can be written as

$$
\Gamma(Z \to \tilde{g}\tilde{g}) = \frac{\alpha M_Z}{48 c_W^2 s_W^2} \left(\frac{\alpha_S}{\pi} \right)^2 \left(\sum_q B^q \right)^2, \tag{20}
$$

with

$$
\begin{aligned}
B^q &= \frac{1}{2}(a_q + v_q \cos 2\theta_q) \left[2C_2(s, m_q, m_q, \tilde{m}_{q_1}) - B_0(s, m_q, m_q) + (m_q^2 - \tilde{m}_{q_1}^2) C_0(s, m_q, m_q, \tilde{m}_{q_1}) \right] \\
&+ \frac{1}{2}(a_q - v_q \cos 2\theta_q) \left[2C_2(s, m_q, m_q, \tilde{m}_{q_2}) - B_0(s, m_q, m_q) + (m_q^2 - \tilde{m}_{q_2}^2) C_0(s, m_q, m_q, \tilde{m}_{q_2}) \right] \\
&- a_q m_q^2 \left[C_0(s, m_q, m_q, \tilde{m}_{q_1}) + C_0(s, m_q, m_q, \tilde{m}_{q_2}) \right] \\
&- \cos 2\theta_q \left[a_{11} C_2(s, \tilde{m}_{q_1}, \tilde{m}_{q_1}, m_q) - a_{22} C_2(s, \tilde{m}_{q_2}, \tilde{m}_{q_2}, m_q) \right] \\
&+ 2 \sin 2\theta_q a_{12} C_2(s, \tilde{m}_{q_1}, \tilde{m}_{q_2}, m_q).
\end{aligned} \tag{21}
$$

In terms of the scalar two and three point functions B_0 and C_0 given in the Appendix, the function C_2 is defined as

$$
\begin{aligned}
C_2(s, m_1, m_2, m_3) &= \frac{1}{2} m_3^2 C_0(s, m_1, m_2, m_3) + \frac{1}{4} + \frac{1}{4} B_0(s, m_1, m_2) - \frac{1}{8s}(m_1^2 + m_2^2 - 2m_3^2) \\
&\times \left[2B_0(s, m_1, m_2) - B_0(0, m_3, m_1) - B_0(0, m_3, m_2) + (2m_3^2 - m_1^2 - m_2^2) C_0(s, m_1, m_2, m_3) \right] \\
&+ \frac{1}{8s}(m_1^2 - m_2^2) \left[B_0(0, m_3, m_1) - B_0(0, m_3, m_2) + (m_2^2 - m_1^2) C_0(s, m_1, m_2, m_3) \right]
\end{aligned} \tag{22}
$$

Note that B_i^q and $\tilde{B}_{ij}^q$ are ultraviolet divergent [the term proportional to the vectorial coupling v_q in B_i^q is finite when one sums over the contributions of the two squarks of a given flavor]. It is only when one sums over a complete isodoublet that the amplitudes are finite. The contributions from (s)bottom loops therefore always have to be included here. On the other hand, if both squarks and quarks of a given generation are degenerate in mass, this generation will not contribute to B_q; see also refs. [10, 11].

3. Numerical Results

We are now in a position to present numerical results for branching ratios of Z and Higgs bosons into pairs of light gluinos. As stated in the Introduction, we will assume minimal (s)particle content, i.e. two Higgs doublets as well as the usual gauge and matter superfields. Moreover, unless stated otherwise we will assume that explicitly SUSY breaking contributions to the squark mass matrices are identical for all flavors. In terms of the parameters introduced in eqs.(2), this implies:

$$m_{\tilde{t}_L} = m_{\tilde{t}_R} = m_{\tilde{b}_R} \equiv m_{\tilde{q}}; \tag{23a}$$

$$A_t = A_b \equiv A \cdot m_{\tilde{q}}. \tag{23b}$$

When computing Higgs branching fractions we have to specify one more parameter, besides those appearing in eqs.(2). We choose this to be the mass m_P of the pseudoscalar Higgs boson. We include full 1–loop corrections [16] from the (s)top and (s)bottom sectors to the mass and mixing angle of the scalar Higgs bosons, including non–logarithmic terms [17].

The values of these parameters are constrained by unsuccessful searches for squarks and Higgs bosons. However, many bounds that have been derived under the assumption that gluinos are heavy are no longer valid if the gluino mass is just a few GeV. In particular, heavy squarks would almost always decay into the corresponding quark plus a gluino. Since a light gluino will lose a substantial fraction of its energy in radiation prior to its decay [3], the missing transverse momentum in such events might be too small for them to pass cuts designed for conventional squark searches at hadron colliders [4]. On the other hand, LEP experiments allow to place bounds on squark masses simply from the measurement of the total and hadronic widths of the Z boson, independent of how the squarks decay. We therefore require all squark mass eigenstates to be heavier than 45 GeV.[3]

Turning to bounds on the Higgs sector, we note that searches for $Z \rightarrow Z^* h$ decays (the Higgs bremsstrahlung process) are little affected by Higgs decay modes as long as it decays hadronically. We have therefore excluded combinations of parameters that violate the m_h–dependent bound on the ZZh coupling derived by the ALEPH collaboration [18].[4] On the other hand, searches for $Z \rightarrow hP$ decays do make use of the assumption that Higgs bosons accessible at LEP decay predominantly into pairs of b quarks or τ leptons; some sort of heavy fermion tagging is necessary in order to suppress the QCD 4–jet background. Thus bounds from searches for associate hP production may not apply if these Higgs bosons have a significant or even dominant branching ratio into light gluinos. We have therefore not included these bounds in our analysis.

In figs. 2a,b we show results for $Br(h \rightarrow \tilde{g}\tilde{g})$ for $m_{\tilde{q}} = 400$ GeV, $\mu = 200$ GeV and several combinations of m_P and tanβ. The curves have been obtained by varying the A parameter in the region $A < 0$ from its minimal allowed value, defined by $m_{\tilde{t}_1} = 45$ GeV, to the point where $m_{\tilde{t}_1}$ is maximized ($A_t = -\mu \cot\beta$). The results are presented as a function of the light

[3]In practice we only include contributions from third generation (s)quarks; large masses for squarks of the first and second generation would not affect our results. This makes it even less likely that the parameter space relevant for us is constrained significantly by sparticle searches at $p\overline{p}$ colliders.

[4]The OPAL collaboration has recently published improved Higgs search limits [19]. Unfortunately the information given in their paper does not allow to extract an upper bound on the ZZh coupling.

stop mass $m_{\tilde{t}_1}$, since it is more directly accessible experimentally than the A parameter.

We find that for $m_P^2 \gg m_Z^2$ the dependence of the branching ratio on either m_P or tanβ is quite weak. In this limit the light Higgs scalar h couples to quarks, leptons and massive gauge bosons with the same strength as the single physical Higgs boson of the SM [5, 6]. If in addition $m_{\tilde{t}_1}^2 \ll m_{\tilde{q}}^2$, the $h\tilde{t}_1\tilde{t}_1$ coupling $(s_2^{11})^q$ is [14] proportional to the off–diagonal element of the stop mass matrix (2a); this implies that $(s_2^{11})^q \propto \left(m_{\tilde{t}_2}^2 - m_{\tilde{t}_1}^2\right)/m_W \propto m_{\tilde{q}}^2/m_W$ in this limit. For fixed $m_{\tilde{t}_1}$, the square of this coupling, and hence $\Gamma(h \to \tilde{g}\tilde{g})$, therefore grows like the fourth power of $m_{\tilde{q}}$. This explains why decays into light gluinos completely dominate all other decay modes for small $m_{\tilde{t}_1}$ and large m_P, as shown in figs. 2. Even for the rather modest value of $m_{\tilde{q}}$ chosen here this new deay mode can suppress the branching ratios into $b\bar{b}$ or $\tau^+\tau^-$ pairs by more than a factor of 1,000.

The situation is more complicated for smaller values of m_P. If m_P is of order m_Z or less, the scalar mixing angle changes such that $\Gamma(h \to b\bar{b}) \propto \tan^2\beta$ for large tanβ, while the coupling of h to top quarks becomes $\propto$ cotβ. Both effects lead to a rapid decrease of $Br(h \to \tilde{g}\tilde{g})$ if m_P is not large. Even for tanβ as small as 2, fig. 2b, for small m_P the gluino mode is important only if $m_{\tilde{t}_1}$ is very small, although it might still remain measurable over a wider region of parameter space. Notice also that the branching ratio now depends on A_t and μ separately, not only on the combination $A_t + \mu\cot\beta$ that appears in the stop mass matrix (dotted and short dashed curves).

Fig. 3 shows results for the decay of the pseudoscalar Higgs boson into light gluinos. The couplings of P are uniquely determined by tanβ, independent of m_P; our result is therefore almost independent of m_P as long as it lies in the region $2m_b \leq m_P \leq 2m_t$.[5] In fig. 3 we took $m_P = 50$ GeV, and chose the same values of $m_{\tilde{q}}$ and μ as in fig. 2.

We see from fig. 3 that the maximal contribution of the $\tilde{g}\tilde{g}$ final state to the total decay width of P is clearly much smaller than in case of the scalar h. The reason is that there is no $P\tilde{t}_1\tilde{t}_1$ coupling [5]. The $P\tilde{t}_1\tilde{t}_2$ coupling (13) can still become quite large if $m_{\tilde{t}_1} \ll m_{\tilde{q}}$; however, this coupling only contributes through a diagram containing one *heavy* stop $\tilde{t}_2$ in the loop, giving a factor $1/m_{\tilde{t}_2}^2$ from the loop integration. As a result, for fixed $m_{\tilde{t}_1}$ the $P \to \tilde{g}\tilde{g}$ partial width becomes independent of $m_{\tilde{q}}$ once $m_{\tilde{q}}^2 \gg m_t^2$, as compared to a growth $\propto m_{\tilde{q}}^4$ in case of the scalar Higgs boson h. Nevertheless, for small tanβ the $\tilde{g}\tilde{g}$ mode can still be quite important even for the pseudoscalar, suppressing the branching ratio into b and τ pairs by a factor of about 3. For larger tanβ the width for these standard modes grows $\propto \tan^2\beta$, while the Ptt and $P\tilde{t}_1\tilde{t}_1$ couplings decrease; the $\tilde{g}\tilde{g}$ mode therefore becomes negligible for tan$\beta > 5$ or so.

The same parameters that determine the partial widths for Higgs boson decays into light gluinos also determine $\Gamma(Z \to \tilde{g}\tilde{g})$; one might therefore hope to place bounds on the former by constraining the latter. However, fig. 4 shows that $\Gamma(Z \to \tilde{g}\tilde{g})$ depends only very weakly on details of $L - R$ squark mixing, described by the parameters A, μ and tanβ. Recall that we assume squarks of the first two generations to be degenerate in mass; these generations do therefore not contribute to the $Z \to \tilde{g}\tilde{g}$ amplitude. Mixing, as well as splitting of mass

[5]We include $P \to Zh$ decays where kinematically allowed when computing the total decay width of P, but this contribution is usually very small.

eigenstates, can be sizable for stop squarks; we have seen above that the partial widths for Higgs decays into light gluinos strongly depend on these details of the stop sector. In contrast, the $Z\tilde{t}_i\tilde{t}_j$ couplings are pure gauge couplings, which are usually modified only weakly by squark mixing. Moreover, reducing $m_{\tilde{t}_1}$ below m_t has little effect on the loop integrals. It is also important to keep in mind that we get a meaningful result only after summing over all four squark mass eigenstates of a given generation; this summation further reduces the dependence on the details of the squark mass matrices. Even increasing $m^2_{\tilde{t}_R}$ and $m^2_{\tilde{b}_R}$ to $2 \cdot m^2_{\tilde{t}_L}$ (dotted curve) reduces $\Gamma(Z \to \tilde{g}\tilde{g})$ by at most 20%, since the couplings of right–handed squarks to the Z are rather weak anyway.

We see that within our assumptions $Br(Z \to \tilde{g}\tilde{g}) \leq 7 \cdot 10^{-4}$. Somewhat larger branching ratios are possible [11] if we allow mass splitting between squarks of the first two generations, but this splitting is tightly constrained by bounds on flavor–changing neutral currents, most notably $K^0\overline{K^0}$ mixing [20]. The contribution of light gluinos to the total or hadronic decay width of the Z boson is thus always well below the present experimental error.[6] One will therefore have to devise cuts that allow to distinguish gluino pairs from pairs of light quarks. One obvious example is to seach for events with missing energy. However, this requires knowledge of the gluino fragmentation function which is not very well understood if gluinos are light [3]. Moreover, there are substantial backgrounds, both instrumental (from the mismeasurement of jet energies) and irreducible (from the production and subsequent semi–leptonic decay of heavy quarks).

A more promising strategy might therefore be [22] to search for $\tilde{g}$ pairs by looking for their decay vertices. Farrar recently concluded [2] that $m_{\tilde{g}} \geq 3$ GeV, mostly due to constraints from searches for radiative Υ decays into $\tilde{g}\tilde{g}$ bound states; the decay vertices of gluinos of this mass should be readily detectable, unless the LSP mass is very close to the gluino mass. If the LSP is an (almost) massless photino, the gluino lifetime is approximately given by [2]

$$\tau_{\tilde{g}} \simeq 5 \cdot 10^{-18} \text{ sec} \cdot \frac{3 \text{ GeV}}{m_{\tilde{g}}} \left(\frac{m_{\tilde{q}}}{m_{\tilde{g}}} \right)^4 . \tag{24}$$

This gives a mean flight path of approximately 0.1 mm (2 cm) for $m_{\tilde{q}} = 50$ (100) GeV and $m_{\tilde{g}} = 3$ GeV. A detailed Monte Carlo study will be necessary to determine whether detection of light gluinos at LEP via these or other [23, 22] methods is feasible.

Finally, in figs. 5 a–c we show scatter plots in the planes spanned by any two of the three branching ratios discussed above. These plots have been obtained by randomly choosing sets of input parameters within the limits 50 GeV $\leq m_{\tilde{q}} \leq$ 500 GeV, 50 GeV $\leq \mu \leq$ 1 TeV, 25 GeV $\leq m_P \leq$ 500 GeV, $1.2 \leq \tan\beta \leq 25$ and $-4 \leq A \leq 4$, excluding combinations of parameters that lead to too light a squark eigenstate or too large a ZZh coupling. Fig. 5a shows a weak positive correlation between the h and P branching ratios: Most points with $Br(P \to \tilde{g}\tilde{g}) > 10\%$ also have $Br(h \to \tilde{g}\tilde{g}) > 1\%$ at least; there are some counter–examples, however. Also, the number of points with $Br(h \to \tilde{g}\tilde{g}) > 10\%$ is clearly much larger than that with $Br(P \to \tilde{g}\tilde{g}) > 10\%$; this is not surprising, given the very different dependence of

[6] Light gluinos will also contribute to the hadronic and total Z boson width through loop diagrams; however these contributions are generally small [21].

the two decay amplitudes on $\tan\beta$ and on $m_{\tilde{q}}$, as discussed above.

Fig. 5b shows that the h and Z branching ratios into light gluinos are anti–correlated. The reason is that, as explained above, a large $Br(h \to \tilde{g}\tilde{g})$ can most easily be obtained if $m^2_{\tilde{q}} \gg m^2_{\tilde{t}_1}$ is sizable, which suppresses $\Gamma(Z \to \tilde{g}\tilde{g})$. Fig. 5c shows no significant correlation between the P and Z branching ratios into gluinos. Taken together, figs. 5 a–c demonstrate that the measurement of any one of the three branching ratios discussed here would not allow one to predict, or even significantly constrain, the other two.

4. Summary and Conclusions

In this paper we computed the decay widths of Higgs and Z bosons into pairs of light gluinos. Although this decay only occurs at the 1–loop level, it can dominate the total decay width of the light scalar Higgs boson by a large factor if the off–diagonal entry of the stop mass matrix is large, in which case the lighter stop eigenstate is expected to be (much) lighter than the other squarks. The contribution of the $\tilde{g}\tilde{g}$ final state to the total decay width of the pseudoscalar Higgs boson can also be quite sizable, although not as large as for the sclar Higgs boson. In contrast, the branching ratio for $Z \to \tilde{g}\tilde{g}$ decays is always below one permille; a dedicated search will be necessary to explore this possibility experimentally. Unfortunately placing bounds on, or a measurement of, $Br(Z \to \tilde{g}\tilde{g})$ will not teach us much about Higgs branching ratios into light gluinos.

What are the consequences of large $Br(h, P \to \tilde{g}\tilde{g})$? As discussed in sec. 3, searches for $Z \to Z^*h$ decays at LEP1, or $e^+e^- \to Zh$ production at LEP2, are probably not much affected, since here it is usually not necessary to tag the Higgs decay products in order to isolate a detectable signal. Some heavy fermion tagging *is* necessary in searches for associate hP production, however. Notice that at present only the combined limits on Z^*h and hP production allow one to place lower bounds on m_h and m_P [18, 19]. We are therefore forced to conclude that the present bounds on the masses of the Higgs bosons of the MSSM may well be invalid if there are indeed gluinos with mass of only a few GeV.

On the other hand, in the long run the existence of light gluinos need not hamper Higgs searches at e^+e^- colliders. While gluino pairs do look different from b and especially τ pairs, they also ought to differ sufficiently from standard light quark or gluon jets to allow the suppression of QCD backgrounds. The impact of light gluinos on Higgs searches at hadron colliders might be less benign, however. The huge background from pure QCD processes means that gluino pairs cannot be used as a Higgs signal here. We have seen that in the presence of light gluinos the branching ratio of scalar Higgs bosons into b or τ pairs might be reduced by a factor of 1,000 or more. The same suppression factor applies for decays into (virtual) gauge bosons. The $h \to ZZ^* \to 4$ leptons signal would then be undetectable; however, within the MSSM it is at best marginal anyway [24].

Potentially most unfortunate for the prospects of MSSM Higgs searches at hadron colliders would be a reduction of the $h \to \gamma\gamma$ signal, which is usually regarded to be the most promising way to search for h [24]. However, the same $h\tilde{t}_1\tilde{t}_1$ coupling that can give rise to a large partial width for the $\tilde{g}\tilde{g}$ final state also contributes to the $h \to \gamma\gamma$ and $h \to gg$ partial widths, via $\tilde{t}_1$ loops. A proper investigation of the possible impact of the existence of light

gluinos on Higgs boson searches at hadron colliders therefore necessitates an analysis of all SUSY loop contributions to Higgs production and decay, which is beyond the scope of this paper. We nevertheless hope that our results will give additional urgency to the effort to either detect or definetely exclude the existence of gluinos with mass of a few GeV.

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Appendix: Complete result

Taking into account the mass of the gluino, the partial decay width of the CP–even and CP–odd Higgs bosons into gluino pairs is given by

$$\Gamma(H_k \to \tilde{g}\tilde{g}) = \frac{1}{2}\alpha M_{H_k} \left(\frac{\alpha_S}{\pi}\right)^2 \beta_{\tilde{g}}^p \left[\sum_q \left(\sum_{i=1}^2 A_k^i + \sum_{i,j=1}^2 \tilde{A}_k^{ij}\right)\right]^2 \tag{A.1}$$

where $\beta_{\tilde{g}} = 1 - 4M_{\tilde{g}}^2/M_{H_k}^2$ is the velocity of the final gluinos with $p = 1$ for $k = 1, 2$ and $p = 3$ for $k = 3$. The amplitudes A_k^i come from the diagrams Fig. 1a and $\tilde{A}_k^{ij}$ come from the diagrams Fig. 1b; they are given by

$$\begin{aligned}
A_{1,2}^i &= (s_{1,2})^q (v_i^2 - a_i^2) \Big[B_0(s, m_q, m_q) + ({m_{\tilde{g}}}^2 + m_q^2 + {m_{\tilde{q}_i}}^2) C_0(s, m_q, m_q, m_{\tilde{q}_i}) + 4{m_{\tilde{g}}}^2 \\
&\qquad C_+(s, m_q, m_q, m_{\tilde{q}_i})\Big] + 2m_q m_{\tilde{g}} (v_i^2 + a_i^2) \left[C_0(s, m_q, m_q, m_{\tilde{q}_i}) + 2C_+(s, m_q, m_q, m_{\tilde{q}_i})\right] \\
A_3^i &= (s_3)^q (v_i^2 - a_i^2) \Big[-B_0(s, m_q, m_q) + ({m_{\tilde{g}}}^2 + m_q^2 - {m_{\tilde{q}_i}}^2) C_0(s, m_q, m_q, m_{\tilde{q}_i})\Big] \\
&\qquad + 2m_q m_{\tilde{g}} (v_i^2 + a_i^2) C_0(s, m_q, m_q, m_{\tilde{q}_i}) \\
\tilde{A}_{1,2}^{ij} &= \left(\tilde{s}_{1,2}^{ij}\right)^q \Big[m_q (v_i v_j - a_i a_j) C_0(s, m_{\tilde{q}_i}, m_{\tilde{q}_j}, m_q) - 2m_{\tilde{g}} (v_i v_j + a_i a_j) C_+(s, m_{\tilde{q}_i}, m_{\tilde{q}_j}, m_q)\Big] \\
\tilde{A}_3^{ij} &= \left(\tilde{s}_3^{ij}\right)^q \Big[m_q (v_i a_j - v_j a_i) C_0(s, m_{\tilde{q}_i}, m_{\tilde{q}_j}, m_q) - 2m_{\tilde{g}} (v_i a_j + a_i v_j) C_-(s, m_{\tilde{q}_i}, m_{\tilde{q}_j}, m_q)\Big]
\end{aligned} \tag{A.2}$$

with the couplings $(s_k)^q$ and $(\tilde{s}_k^{ij})^q$ given in eqs. (6–10) and

$$v_1 = \frac{1}{2}(\cos\theta_q - \sin\theta_q) = a_2 \quad , \quad a_1 = \frac{1}{2}(\cos\theta_q + \sin\theta_q) = -v_2 \tag{A.3}$$

The functions $C_+(s, m_1, m_2, m_3)$ is defined as

$$\begin{aligned}
C_+(s, m_1, m_2, m_3) &= \frac{1}{2s\beta_{\tilde{g}}^2} \Big[2B_0(s, m_1, m_2) - B_0({m_{\tilde{g}}}^2, m_3, m_1) - B_0({m_{\tilde{g}}}^2, m_3, m_2) \\
&\qquad + (2{m_{\tilde{g}}}^2 + 2m_3^2 - m_1^2 - m_2^2) C_0(s, m_1, m_2, m_3)\Big] \\
C_-(s, m_1, m_2, m_3) &= \frac{1}{2s} \Big[B_0({m_{\tilde{g}}}^2, m_3, m_2) - B_0({m_{\tilde{g}}}^2, m_3, m_1) + (m_1^2 - m_2)^2 C_0(s, m_1, m_2, m_3)\Big]
\end{aligned} \tag{A.4}$$

with the scalar two and three point functions, B_0 and C_0 defined as

$$\begin{aligned}
B_0(s, m_1, m_2) &= \frac{(2\pi\mu)^{n-4}}{i\pi^2} \int \frac{d^n k}{(k^2 - m_1^2 + i\epsilon)[(k-q)^2 - m_2^2 + i\epsilon]}, \\
C_0(s, m_1, m_2, m_3) &= \frac{(2\pi\mu)^{n-4}}{i\pi^2} \int \frac{d^n k}{[(k-p_1)^2 - m_1^2 + i\epsilon][(k-p_2)^2 - m_2^2 + i\epsilon](k^2 - m_3^2 + i\epsilon)}.
\end{aligned} \tag{A.5}$$

where n is the space–time dimension and μ the renormalisation scale. After integration over the internal momentum k, the function B_0 is given by [γ_E is Euler's constant]

$$B_0(s, m_1, m_2) = \frac{1}{\epsilon} - \gamma_E + + \log\frac{4\pi\mu^2}{m_1 m_2} + 2 + \frac{m_1^2 - m_2^2}{2s} \log\frac{m_2^2}{m_1^2} + \frac{x_+ - x_-}{4s} \log\frac{x_-}{x_+}$$

with

$$x_\pm = s - m_1^2 - m_2^2 \pm \sqrt{s^2 - 2s(m_1^2 + m_2^2) + (m_1^2 - m_2^2)^2} \tag{A.6}$$

The three point scalar function C_0 for $p_1^2 = p_2^2 = m_{\tilde{g}}^2$ and $s = (p_1 + p_2)^2$ is given in integral form by

$$C_0(s, m_{\tilde{g}}, m_1, m_2, m_3) = -\int_0^1 dy \int_0^y dx \left[ay^2 + bx^2 + cxy + dy + ex + f\right]^{-1}, \tag{A.7}$$

where

$$a = m_{\tilde{g}}{}^2 \ , \quad b = s \ , \quad c = -s \ , \quad d = m_2^2 - m_3^2 - m_{\tilde{g}}{}^2 \ , \quad e = m_1^2 - m_2^2 \ , \quad f = m_3^2 - i\epsilon. \tag{A.8}$$

C_0 can be expressed in terms of a sum of Spence functions $\mathrm{Li}_2(x) = -\int_0^1 dt \log(1 - xt)/t$:

$$C_0(s, m_{\tilde{g}}, m_1, m_2, m_3) = -\frac{1}{s\beta_{\tilde{g}}} \sum_{i=1}^{3} \sum_{j=+,-} (-1)^i \left[\mathrm{Li}_2\left(\frac{x_i}{x_i - y_{ij}}\right) - \mathrm{Li}_2\left(\frac{x_i - 1}{x_i - y_{ij}}\right)\right], \tag{A.9}$$

where we have defined

$$\begin{aligned}
x_1 &= \frac{2d + e(1-\beta_q)}{2s\beta_q} + \frac{1}{2}(1-\beta_q) \ , & y_{1\pm} &= \frac{-c - e \pm \sqrt{(c+e)^2 - 4b(a+d+f)}}{2b}, \\
x_2 &= \frac{2d + e(1-\beta_q)}{s\beta_q(1+\beta_q)} \ , & y_{2\pm} &= \frac{-d - e \pm \sqrt{(d+e)^2 - 4f(a+b+c)}}{2(a+b+c)}, \\
x_3 &= -\frac{2d + e(1-\beta_q)}{s\beta_q(1-\beta_q)} \ , & y_{3\pm} &= \frac{-d \pm \sqrt{d^2 - 4af}}{2a}.
\end{aligned} \tag{A.10}$$

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Figure Captions

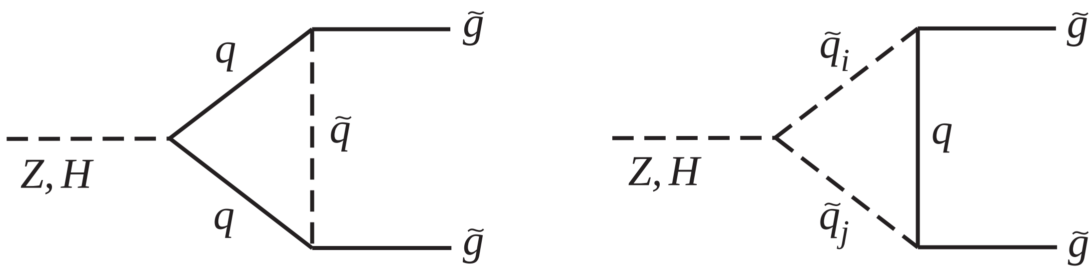
Fig.1 The decays $Z, h, P \to \tilde{g}\tilde{g}$ proceed via 1–loop diagrams with two quark propagators and one squark propagator (1a), or two (possibly different) squark propagators and one quark propagator (1b).

Fig.2 The branching ratio for the decay of the light scalar Higgs boson h of the MSSM into a pair of light gluinos is shown as a function of the mass of the lighter stop eigenstate $\tilde{t}_1$, for the ratio of vacuum expectation values tan$\beta = 25$ (2a) and 2 (2b), respectively. We have fixed $m_{\tilde{q}} = 400$ GeV, and $\mu = 200$ GeV for all cases except the dotted curve in fig. 2b, which is for $\mu = 500$ GeV. The curves have been obtained by varying the A parameter in the region $A < 0$, as explained in the text.

Fig.3 The branching ratio for the decay of the pseudoscalar Higgs boson P of the MSSM into a pair of light gluinos is shown as a function of $m_{\tilde{t}_1}$, for $m_{\tilde{q}} = 400$ GeV, $\mu = 200$ GeV, $m_P = 50$ GeV and three different values of tanβ as indicated. The curves for small tanβ terminate well below the possible maximum of $m_{\tilde{t}_1}$ since here substantial 1–loop corrections to the scalar Higgs sector are necessary to evade the ALEPH bound [18] on the ZZh coupling; this leads to a lower bound on $|A_t|$ in these cases.

Fig.4 The branching ratio for the decay of the Z boson into a pair of light gluinos is shown as a function of the mass $m_{\tilde{q}_L}$ of $SU(2)$ doublet squarks, for $\mu = 200$ GeV and various combinations of A and tanβ. The solid and dashed curves have been derived using our usual assumption of equal explicit SUSY breaking masses for $SU(2)$ doublet and singlet squarks, while the dotted curve is for $m^2_{\tilde{q}_R} = 2m^2_{\tilde{q}_L}$.

Fig.5 Scatter plots in the planes spanned by the branching ratios of h and P (5a), h and Z (5b) and P and Z (5c) into pairs of light gluinos. These plots have been obtained by randomly choosing combinations of input parameters within the boundaries 50 GeV $\leq m_{\tilde{q}} \leq 500$ GeV, 50 GeV $\leq \mu \leq 1$ TeV, 25 GeV $\leq m_P \leq 500$ GeV, $1.2 \leq$ tan$\beta \leq 25$ and $-4 \leq A \leq 4$. Combinations of parameters that violate the bounds on squark masses and the ZZh coupling discussed in the text have been discarded.


$\tilde{g}$
q
$\tilde{q}$
Z, H
q
$\tilde{g}$
$\tilde{g}$
$\tilde{q}_i$
q
Z, H
$\tilde{q}_j$
$\tilde{g}$

(a) (b)

This figure "fig1-2.png" is available in "png" format from:

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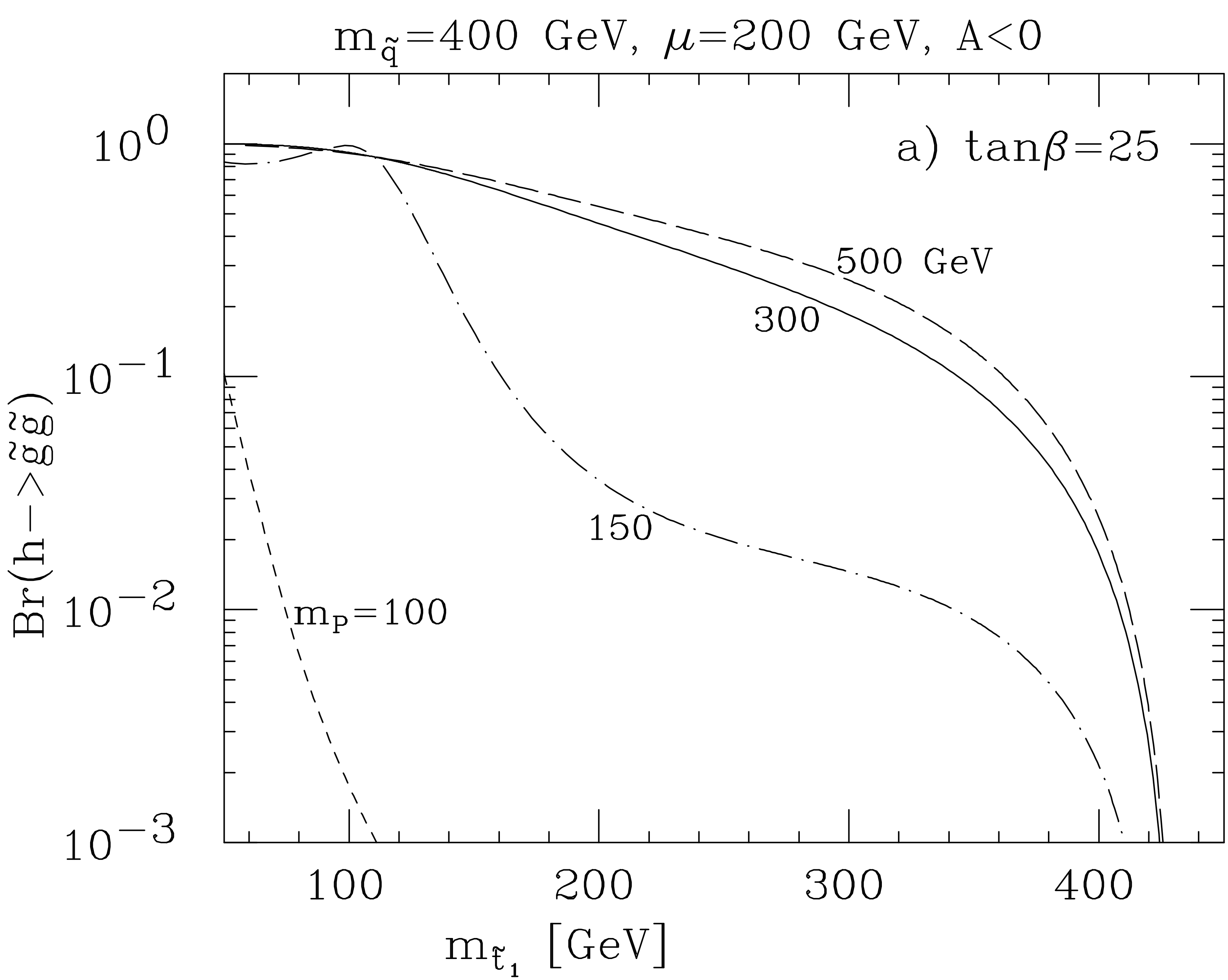
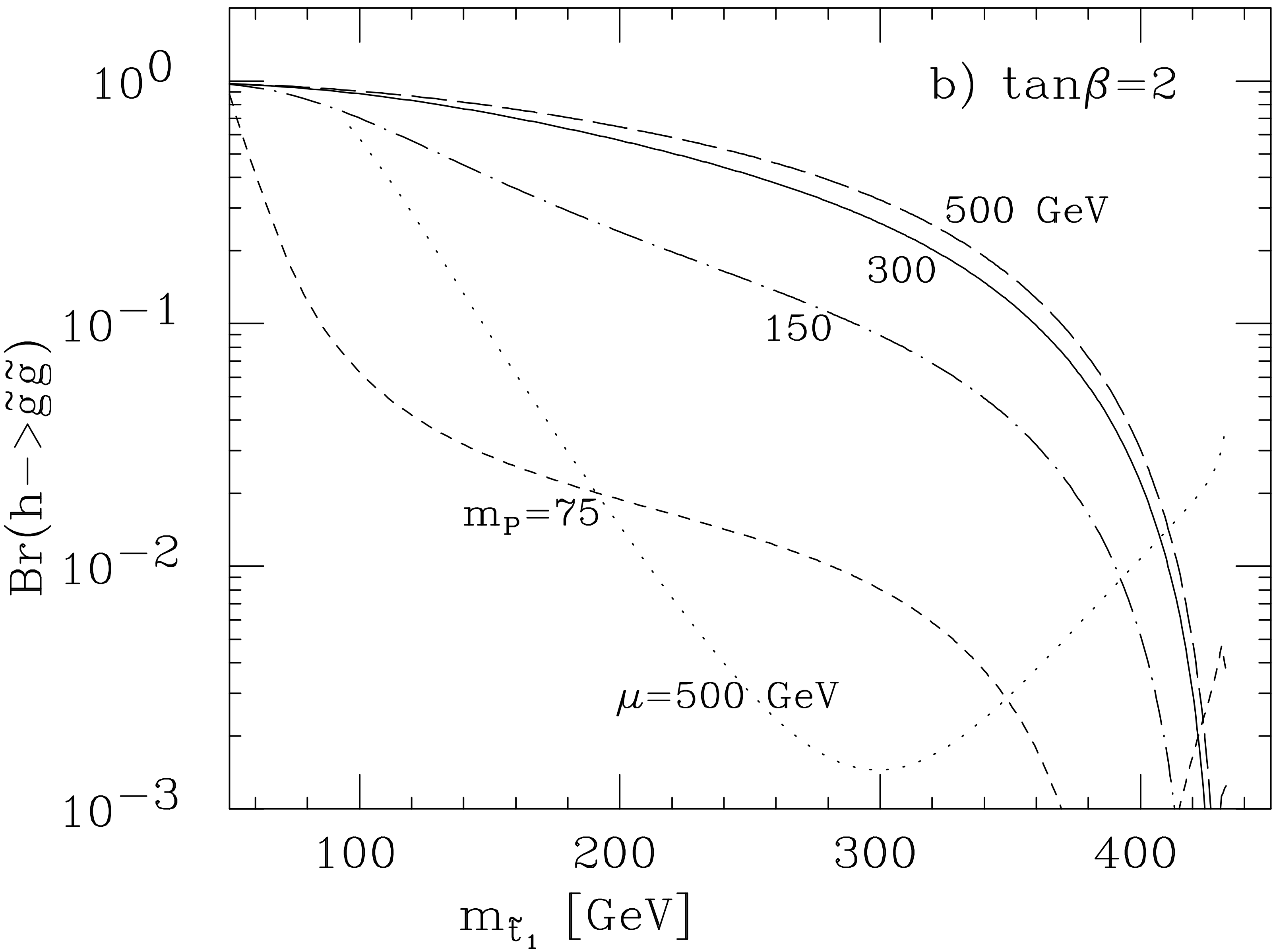
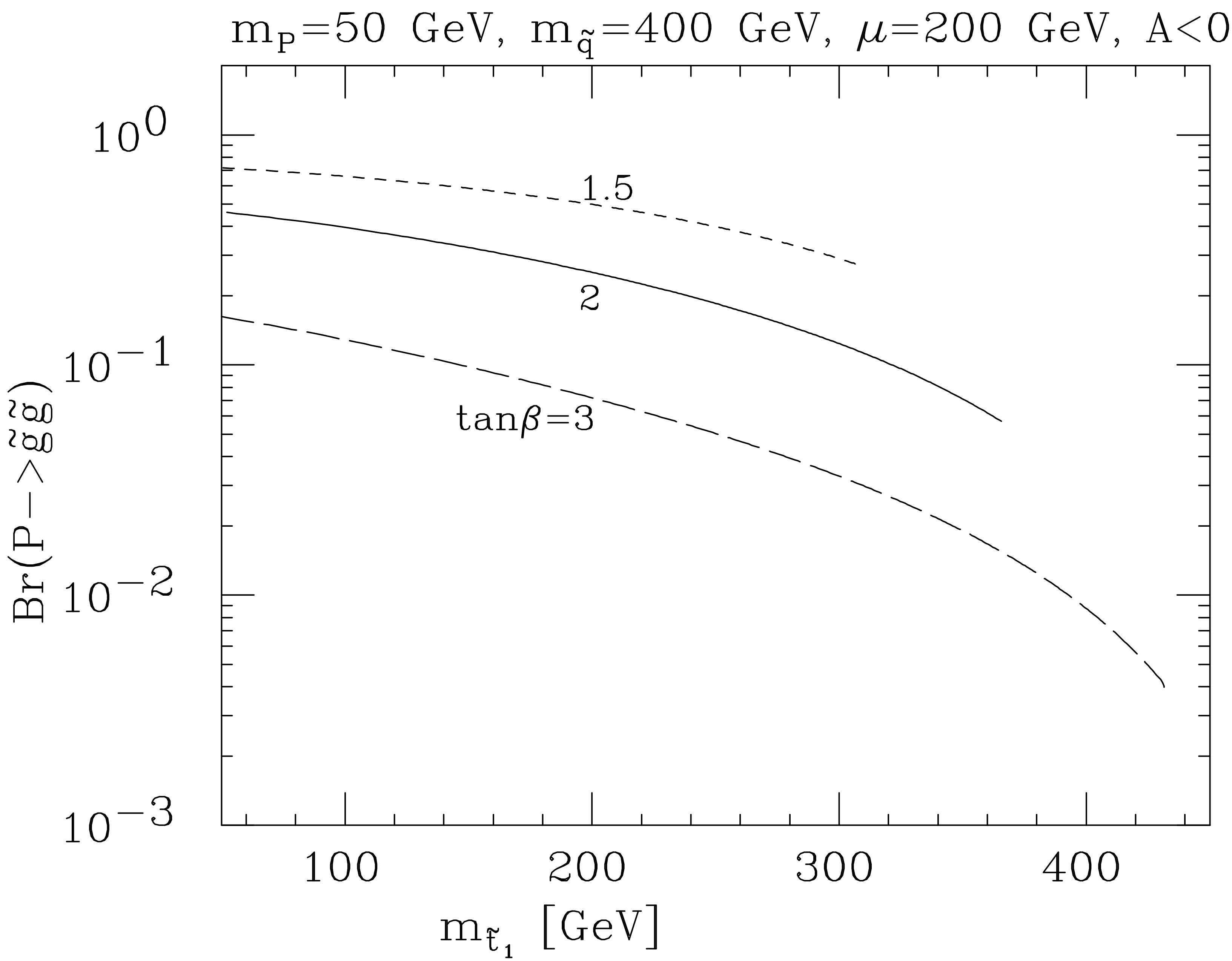
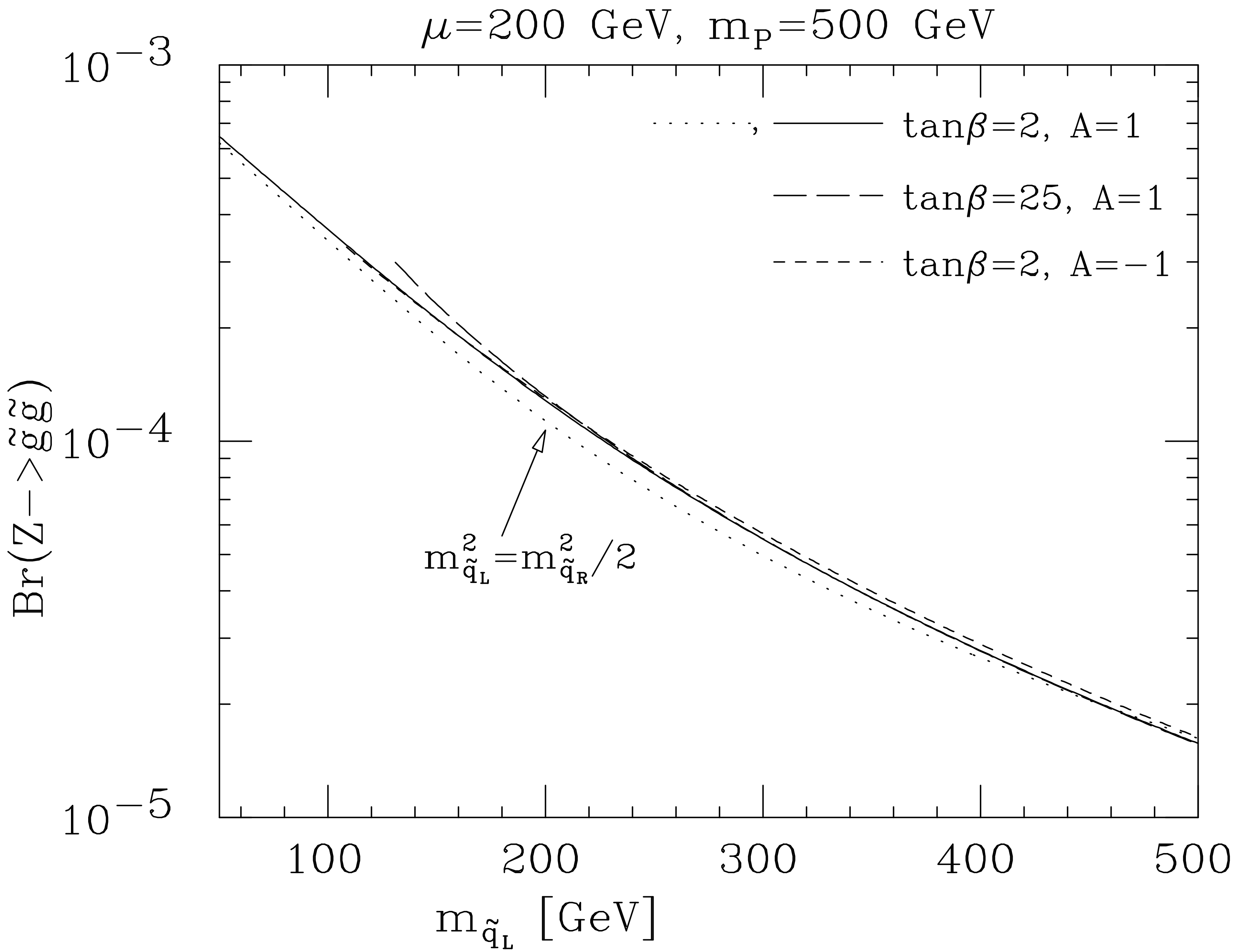
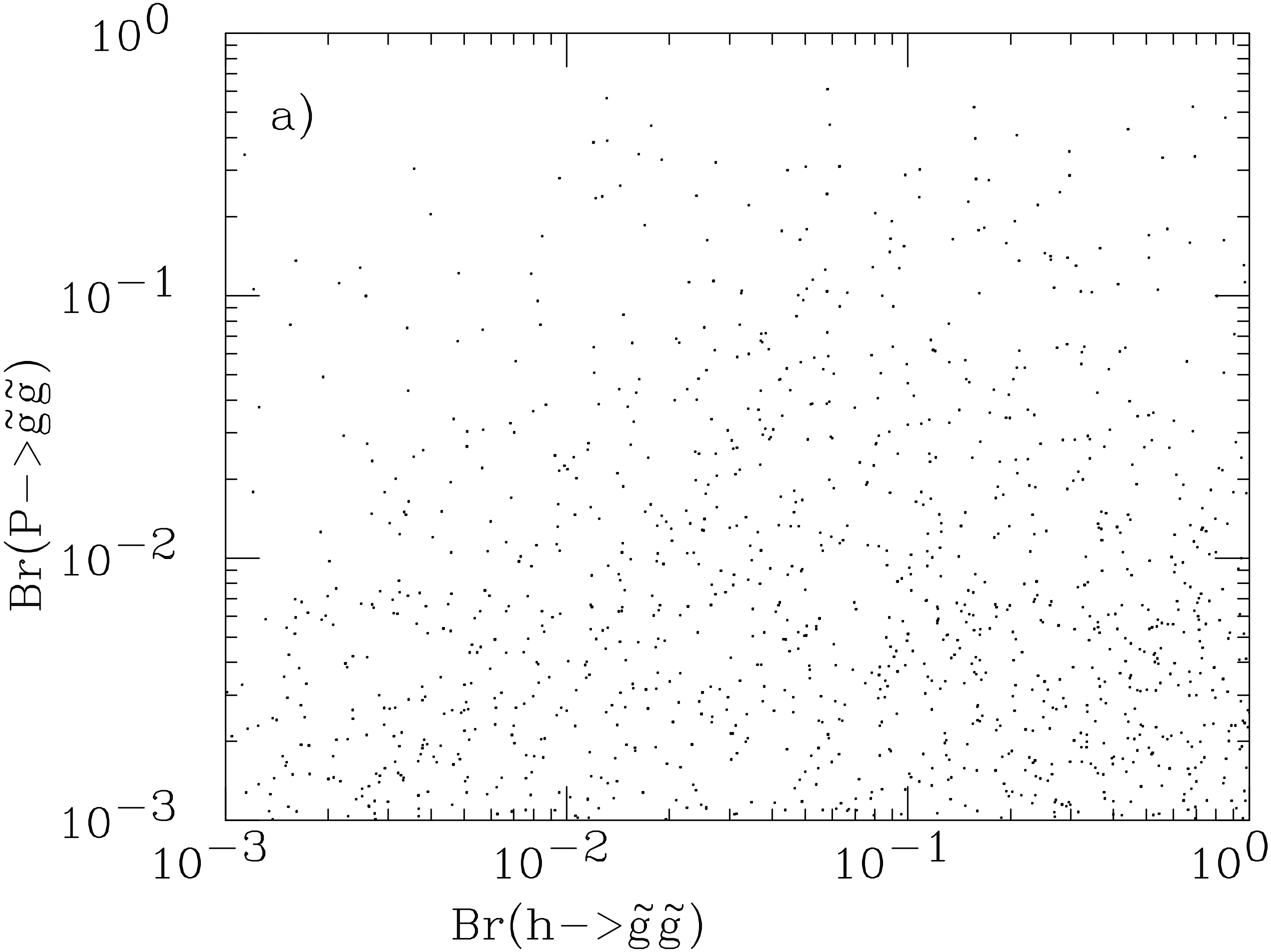


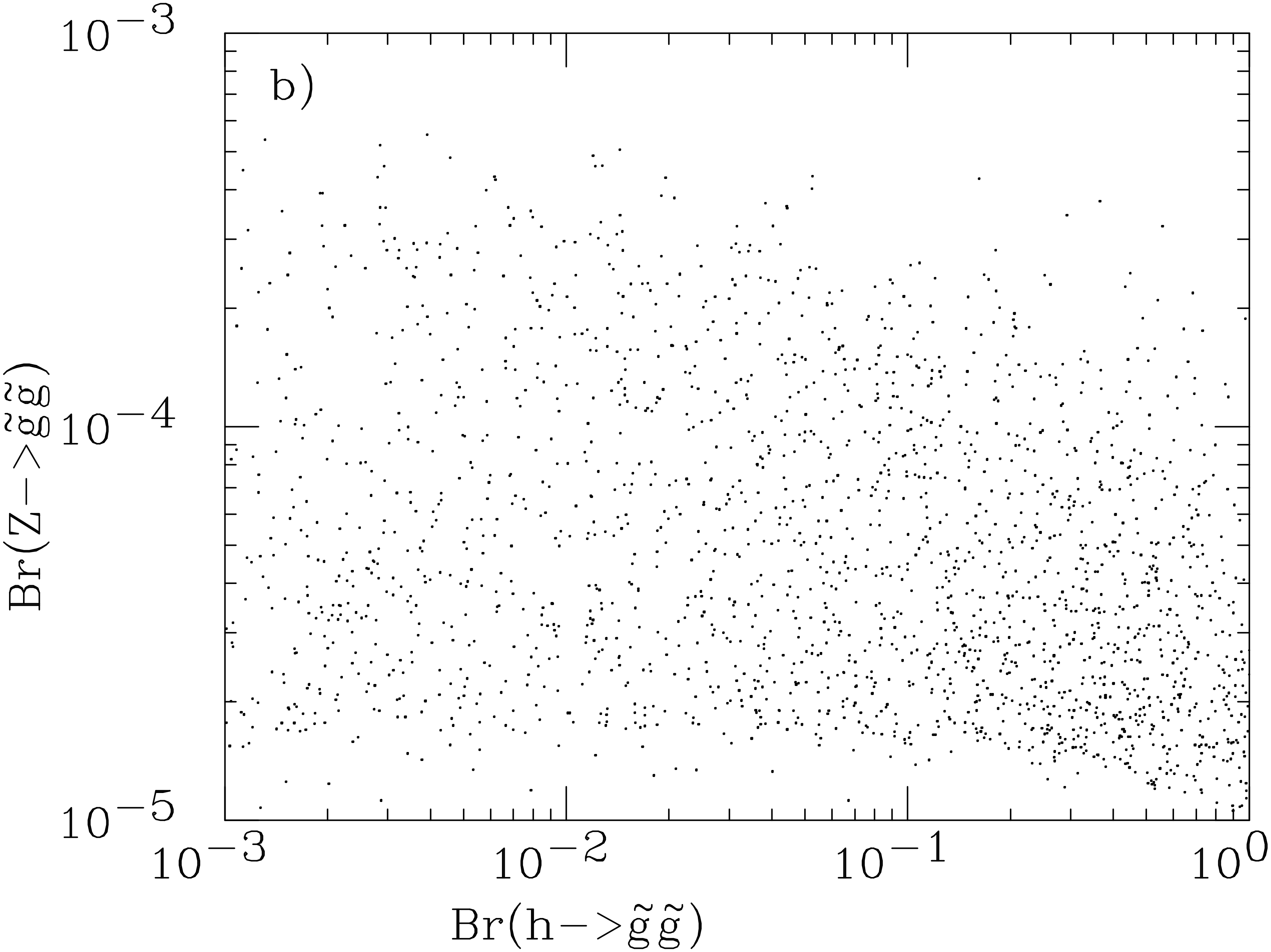
$m_{\tilde{q}}$=400 GeV, μ=200 GeV, A<0
a) tanβ=25
500 GeV
300
150
m_P=100
10^0
10^{-1}
10^{-2}
10^{-3}
Br(h−>$\tilde{g}\tilde{g}$)
100
200
300
400
$m_{\tilde{t}_1}$ [GeV]

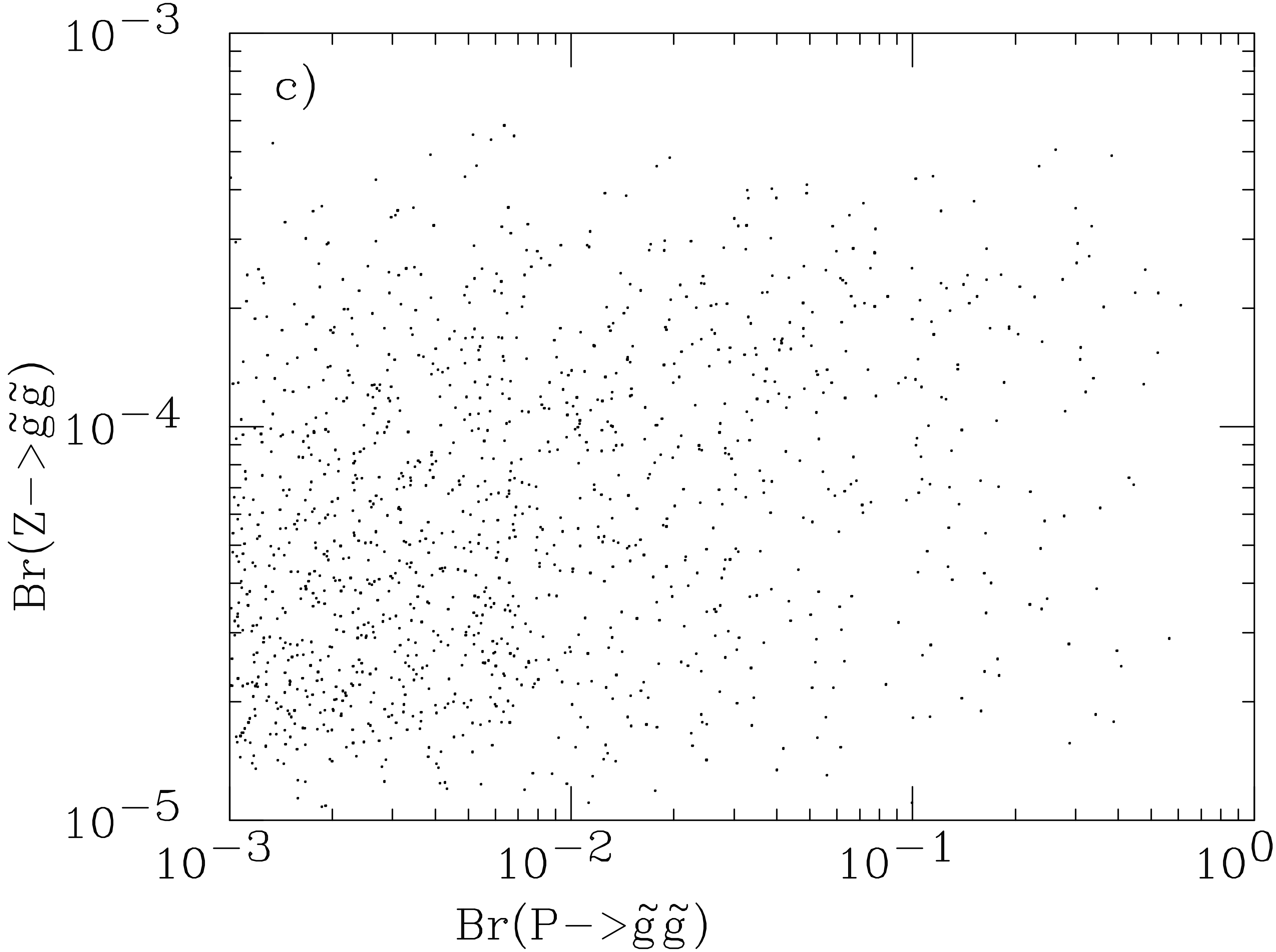

b) tanβ=2
500 GeV
300
150
m_P=75
μ=500 GeV
Br(h->g̃g̃)
10^0
10^-1
10^-2
10^-3
100
200
300
400
m_t̃1 [GeV]


m_P=50 GeV, $m_{\tilde{q}}$=400 GeV, μ=200 GeV, A<0
$Br(P->\tilde{g}\tilde{g})$
10^0
10^{-1}
10^{-2}
10^{-3}
100
200
300
400
$m_{\tilde{t}_1}$ [GeV]
1.5
2
tanβ=3


μ=200 GeV, m_P=500 GeV
10^{-3}
10^{-4}
10^{-5}
Br(Z−>$\tilde{g}\tilde{g}$)
, tanβ=2, A=1
tanβ=25, A=1
tanβ=2, A=−1
$m^2_{\tilde{q}_L}=m^2_{\tilde{q}_R}/2$
100
200
300
400
500
$m_{\tilde{q}_L}$ [GeV]


a)
Br(P−>g̃g̃)
Br(h−>g̃g̃)
10^0
10^−1
10^−2
10^−3
10^−3
10^−2
10^−1
10^0


b)
Br(Z−>g̃g̃)
10−3
10−4
10−5
10−3
10−2
10−1
100
Br(h−>g̃g̃)


c)
10^{-3}
10^{-4}
10^{-5}
$Br(Z->\tilde{g}\tilde{g})$
10^{-3}
10^{-2}
10^{-1}
10^{0}
$Br(P->\tilde{g}\tilde{g})$